MATTER-ANTIMATTER ASYMMETRY AND BARYOGENESIS

By

ASWIN RAJEEV

DU2019MSC0068

DEPARTMENT OF PHYSICS

A THESIS
SUBMITTED
IN PARTIAL FULFILLMENT OF THE REQUIREMENT
FOR THE DEGREE OF

MASTER OF SCIENCE

IN PHYSICS

To

ASSAM DON BOSCO UNIVERSITY

TAPESIA GARDENS, TAPESIA,

SONAPUR – 782 402

ASSAM – INDIA

JULY, 2021

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CERTIFICATE

This is to certify that this report titled

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is a bonafide record of the project work done by **Mr. Aswin Rajeev** (Student ID DU2019MSC0068) at the Department of Physics, School of Fundamental & Applied Sciences, Assam Don Bosco University, under the guidance of **Dr. Debajyoti Dutta**, in partial fulfillment of the requirements of the *PSPR6012: Project Work* of the M.Sc. (Physics) curriculum.

Project Guide

Dr. Debajyoti Dutta

Place: ADBU, Tapesia

Date: July, 2021

Department of Physics
School of Fundamental & Applied Sciences
Assam Don Bosco University
Tapesia, Sonapur – 782 402

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Head of the Department

Dr. Ngangom Aomoa

Place: ADBU, Tapesia

Date: July, 2021

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School of Fundamental & Applied Sciences
Assam Don Bosco University
Tapesia, Sonapur – 782 402

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Director, SFAS

Dr. Monmoyuri Baruah

Place: ADBU, Tapesia

Date: July, 2021

Department of Physics
School of Fundamental & Applied Sciences
Assam Don Bosco University
Tapesia, Sonapur – 782 402

EXAMINATION CERTIFICATE

This is to certify that **Aswin Rajeev** bearing **DU2019MSC0068** of **Department of Physics** has carried out the Project Work in a manner satisfactory to warrant its acceptance and also defended it successfully.

I wish him all success in future endeavours.

Examiners:

1. External Examiner

2. Internal Examiner

3. Internal Examiner

DECLARATION

I, Aswin Rajeev, hereby declare that this report of the project work entitled "**MATTER-ANTIMATTER ASYMMETRY AND BARYOGENESIS** " is the record of the work carried out by me under the supervision of Dr. Debajyoti Dutta, Assam Don Bosco University. The content of this project did not form the basis for the award of any previous degree to me or to the best of my knowledge to anybody else. I also declare that the project has not been submitted by me for any research degree in any other university or institute.

This is being submitted to the Assam Don Bosco University for the degree of Master of Science in Physics.

Aswin Rajeev

Department of Physics

Assam Don Bosco University

Tapesia – 782 402, Assam

India

Place: ADBU, Tapesia

Date: July, 2021

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(Aswin Rajeev)
July, 2021
Department of Physics
Assam Don Bosco University.

ABSTRACT

KEY WORDS: Matter-Antimatter Asymmetry,Baryogenesis,Neutrino Masses

In the world of physics, there is no known explanation for the matter–antimatter asymmetry problem, either in the standard model of particle physics or in the theory of general relativity, and it is a reasonable assumption that the cosmos is neutral with all preserved charges. There are many competing theories to explain the matter-antimatter imbalance that led to Baryogenesis. There is, however, no universally accepted explanation to explain the phenomena. This dissertation is based on current research in the relevant area. The various mechanisms of Baryogenesis and its conditions have been briefly explored.Additionally, studies of CP violations under different circumstances are included. Many Leptogenesis mechanisms have the appealing characteristic of linking observable baryon asymmetry to the creation of neutrino masses and mixing. Computation of baryon to anti-baryon and lepton to anti-lepton ratios using currently accepted cosmological and particle physics data is also done.

Contents

List of Figures

CHAPTER 1

INTRODUCTION

Contents

1.1 An overview

The early stages of the observable Universe are thought to have been preceded by cosmic inflation. This paradigm has the potential to set the stage for a big-bang event. The observed cosmos differs from our expectations in some ways. We don't see any antimatter in our environment. This raises the question of why there is more matter than antimatter. Matter and antimatter are processed in the same way according to the standard model. As a result, baryogenesis is both a cosmological and a particle physics conundrum, and it combines particle physics with cosmology in an intriguing way. What we know about the total baryon density of the cosmos comes from Cosmic Microwave Background observations, which, when combined with the other measurements inside ΛCDM, gives us an accurate estimate of baryon density, $\Omega_B h^2$=0.0223$\pm$ 0.0002. (The h parameter is equal to Hubble's value of 0.697.) From these measurements, it is evident that regular matter accounts for just around 5 percent of the universe's energy density; unfortunately, we don't know why this is the case. This is when the term "baryogensis" comes into play.

Apart from that, primordial nucleosynthesis/Big bang nucleosynthesis provides a second independent determination of η. Light elements were generated during the transition from a hot to a cold cosmos, and their abundance derived from primordial nucleosythesis fits the baryon density we get from the CMB. The free parameter when computing the light element abundance from BBM is baryon density. It turns out that the observed results match the value obtained from the CMB, giving us a coherent

picture of the universe's baryon density. These and other cosmic facts, taken together, give us confidence in the conventional cosmological model of inflation and reheating, as well as the ΛCDM with $\eta \simeq 6 \times 10^{-10}$. Basically, baryogenesis is any process that occurs in the early cosmos that results in a baryon asymmery, and it is expected to happen either during or after warming, resulting in a tiny number of matter over antimatter.Because there is no known mechanism for baryogenesis inside the standard model, it is a significant driver for new physics beyond the standard model.

1.1.1 Consideration of Nucleosynthesis and baryon number

The baryon number density varies to scales like a^3 during the evolution of the universe, where 'a' is a cosmological scale factor. The baryon assymmetry can be defined as follows:

$$\eta \equiv \frac{n_B}{n_\gamma} \tag{1.1}$$

where n_B is the difference in the number of baryons and antibaryons, and n_γ is the photon number density at a particular temperature. T. If there isn't a procedure, The baryon number is conserved at energy scales of ≈1 MeV, resulting in entropy to change the photon number.

$$n_B = \frac{\rho_B}{m_B} = \frac{\Omega_B}{m_b}\rho_c \tag{1.2}$$

where ρ_B is the baryonic energy density. critical density ρ_c= 1.88× 10^{-29} h^2gr cm^{-3} where Hubble parameter H_0 being parametraized by 0.5≤ h ≤ 0.9. $h \equiv H/100$ Km Mpc^{-1} sec^{-1}

$$n_B = 1.1 \times 10^{-5} h^2 \Omega_b cm^{-3} \tag{1.3}$$

present temperature of the background radiation is given as $T_0 = 2.735_0 K$ and rising to

$$n_\gamma \simeq 415 \left(\frac{T_0}{2.735_0 K} \right)^3 cm^{-3} \tag{1.4}$$

Putting equation(1.3) and (1.4) together, we get

$$\eta = 2.65 \times 10^{-8} \Omega_B h^2 (\frac{T_0}{2.735_0 K})^{-3} \tag{1.5}$$

the value of range of η is consistent with deuterium and He_3 primordial abundances is

$$4(3) \times 10^{-10} \leq \eta \leq 7(10) \times 10^{-10} \tag{1.6}$$

the range for $\Omega_B h^2$ can be written as

$$0.015(0.011) \leq \Omega_B h^2 \leq 0.026(0.0038). \tag{1.7}$$

baryon asymmetry can be described in terms of $B \equiv \frac{n_B}{s}$, where s is the entropy density of the universe at temperature T. Hence the equation translates into

$$5.7(4.3) \times 10^{-11} \leq B \leq 9.9(14) \times 10^{-11} \tag{1.8}$$

now the question arises as, with the use of standard cosmological model whether we are able to explain the tiny value of η.

Let us consider initially η=0 and b be the final number density of nucleons that are left after the annihilation have frozen out. At temparature $T \leq 1$ GeV the equilibrium abundance of nucleons and anti nucleons is

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{b*}}{n_\gamma} \simeq \left(\frac{m_p}{T}\right)^{3/2} e^{-m_p/T} \tag{1.9}$$

we know as the universe cool off, the number of nucleons and antinucleons decreases as long as the annihilation rate $\Gamma_{ann} \simeq n_b \langle \sigma_A v \rangle$ is larger than $H \simeq 1.66 g_*^{1/2} T^2 / M_p$ which is the expansion rate of the universe. At $T \simeq 20 MeV$, values of both annihilation rate and expansion rate become close and annihilation freeze out, so that the nucleons and aninucleons cannot annihilate anymore. Hence we get equation 1.9 as

$$\frac{n_b}{n_\gamma} \simeq \frac{n_{b*}}{n_\gamma} \simeq 10^{-18} \tag{1.10}$$

This value is lower than what nucleosynthesis predicts.

Another hypothesis is that the modest value of η can be explained by statistical fluc-

tuations in baryon and antibaryon distributions. At the moment, our galaxy contains around 10^{79} photons. There are around 10^{69} baryons in the comoving volume V that includes our galaxy today, but there are 10^{79} baryons and antibaryons at $T > 1GeV$. As a result of pure statistical fluctuations, an asymmetry

$$(n_b - n_{\bar{b}})/n_b \leq 10^{-39.5}$$

is expected, which is too modest to explain the observed baryon symmetry.

CHAPTER 2

BARYOGENESIS

Contents

2.1 The Three basic conditions for Baryogenesis

As we know, the early universe's process of baryon and antibaryon creation resulted in an asymmetry between matter and antimatter, resulting in the presently heavily matter-dominated cosmos. For the process to be carried out, there are some necessary conditions to be satisfied. The conditions are known as Sakharov's criteria and they are explained below.

•Baryon number(B) violation

•Violation of particle-antiparticle (C) symmetry and parity (CP) symmetry

•Departure from Thermal equilibrium

2.1.1 Baryon number (B) violation

The condition is such that ,we start from a baryon symmetric universe.(B=0 and evolve to a universe where $B \neq 0$) In simple words universe go from no baryons to non zero number of baryons.Baryon number violation interactions are therefore mandatory.

2.1.2 C and CP Violation

Charge Conjugation symmetry (C) and the product of charge conjugation and parity(CP) are not exact symmetries.If they were exact symmetries, we can prove that rate of any process that produces an excess of baryons is equal to the process which produces an excess of antibaryons , which in turn leave no baryon number to be created. ultimately we can say that the thermal average of the baryon number operator B, which is odd under both C and CP, is zero unless those discrete symmetries are violated.

2.1.3 Departure from Thermal Equilibrium

We know for a fact that there would be no preferred direction for time to be defined if the particles in the universe remained in equilibrium, and that the occurrence of any surplus baryons would be precluded by CPT invariance. The presence of superheavy decaying particles in the expanding cosmos satisfies the third criterion of Sakharov, which is known as the out of equilibrium decay process.

Suppose X be a certain candidate with mass m_X in thermal equilibrium at a temperature $T << m_X$, number density is given by

$$n_X \simeq g_X (m_x T)^{3/2} e^{\frac{-m_X}{T} + \frac{\mu_x}{T}}$$

where μ_X is the chemical potential.

We know candidate X is in chemical equilibrium because the change in particle number owing to inelastic scatterings $(X + A \to B + C)$ has a Γ_{inel} rate that is greater than the expansion rate. As a result, one of us can write.

$$\mu_X + \mu_A = \mu_B + \mu_C \tag{2.1}$$

similarly the number density of the antiparticle $\bar{X}$is

$$n_{\bar{X}} \simeq g_X (m_x T)^{3/2} e^{\frac{-m_X}{T} - \frac{\mu_x}{T}}$$

because of the process above mentioned we can write $\mu_{\bar{X}}{=}{-}\mu_X$,

$$\bar{X}X \to \gamma\gamma \tag{2.2}$$

$$B \propto n_X - n_{\bar{X}} = 2g_X(m_X T)^3/2$$

$$e^{\frac{-m_X}{T}} sinh\left(\frac{mu_X}{T}\right) \tag{2.3}$$

if the particles X carries baryon number , B will get a contribution as given above. Also it is crucial that , if X and $\bar{X}$ undergo baryon number violating reactions required by the first condition of sakharov,

$$XX \to \bar{X}X \tag{2.4}$$

then one can conclude $\mu_X = 0$ and the contribution provided by the X particles to the net baryon number disappears.

CHAPTER 3

THE STANDARD MODEL OF PARTICEL PHYSICS

Contents

3.1 Symmetry breaking in SM

The standard model is spontaneously broken non-abelian gauge theory based on the symmetry group $SU(3) \times SU(2) \times U(1)$. One can distinguish between the hypercharge -U(1) gauge group which is brocken and the electromagnetic U(1) gauge group which is unbrocken.

The breaking is done by introducing a hypercharge one, SU(2) complex doublet scalar field.

$$\Phi = \begin{pmatrix} \Phi^+ \\ \Phi^0 \end{pmatrix} \tag{3.1}$$

the $SU(2) \times U(1)_Y$ covarient derivative acting on Φ is given by

$$D_\mu \Phi_i = \partial_\mu \Phi_i + W_\mu^\kappa (L^\kappa)_i^j \Phi_j + B_\mu (L^4)_i^j \Phi_j \tag{3.2}$$

where i,j = 1,2 ; $\kappa = 1, 2, 3$ and

$$(L^\kappa)_i^j = \frac{1}{2} ig(\tau^\kappa)_i^j, (L^4)_i^j = ig' Y \delta_i^j \tag{3.3}$$

where W_μ^κ is the SU(2) gauge fields, g and g' respective gauge couplings, B_μ hypercharge U(1) gauge field and the group generators ($L^a = ig_a T^a$, where a= 1,...4 $g_{1,2,3} = g$ and $g_4 = g'$) The τ^κ are the pauli's matrices and the hypercharge operator is normalised ; $Y\Phi = +\frac{1}{2}\Phi$.

we can sstart by looking at the bosonic sector of SM. The scalar potential constrained by the renormalizibilty and $SU(2) \times U(1)_Y$ invariencec.

$$V(\Phi, \Phi^\dagger) = -m^2(\Phi^\dagger\Phi) + \lambda(\Phi^\dagger\Phi)^2 \tag{3.4}$$

where m^2 and λ are real and positive parameters. A local minimum can be obtained by minimising the scalar potential at $\Phi^\dagger\Phi = \frac{m^2}{2\lambda}$. Hence we can write the vacuum expectation value as

$$\nu \equiv \langle\Phi\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ \upsilon \end{pmatrix} \tag{3.5}$$

also,

$$M_{ab}^2 = 2g_a g_b Re(\nu^\dagger\tau^a\tau^b\nu) = g_a g_b \nu^\dagger(\tau^a\tau^b + \tau^b\tau^a)\nu, \tag{3.6}$$

using the above equation , mass matrix of gauge boson can be computed. Along with that it is important note that $\tau_i\tau_j + \tau_j\tau_i = 2\delta^{ij}$ and $\nu^\dagger\nu = \upsilon^2/2$. With the help of these one can get twp mass degenerate states W^1 and W^2 with $m_w^2 = \frac{1}{2}g^2\upsilon^2$ and a non-diagonal matrix which in the $w^3 - B$ basis is given by

$$M^2 = \frac{\upsilon^2}{4}\begin{pmatrix} g^2 & -gg' \\ -gg' & g'^2 \end{pmatrix} \tag{3.7}$$

We know this matrix can be easily diaogonalized by

$$O = \begin{pmatrix} cos\theta_W & -sin\theta_W \\ sin\theta_W & cos\theta_W \end{pmatrix} \tag{3.8}$$

where the Weinberg angle is given as

$$sin\theta_W = \frac{g'}{\sqrt{g^2 + g'^2}} \tag{3.9}$$

Let us define a new basis for the lie algebra

$$\widetilde{L} \equiv O_{ab} L_b \tag{3.10}$$

from the above equation we can deduce $\tilde{L}_\kappa = L_\kappa$ for k=1,2 and

$$\widetilde{L}_3 = L_3 cos\theta_W - L_4 sin\theta_W = \frac{ig}{cos\theta_w}[\tau^3 - Qsin^2\theta_w] \tag{3.11}$$

$$\widetilde{L}_4 = L_3 sin\theta_W + L_4 cos\theta_W = ieQ \tag{3.12}$$

where

$$e = gsin\theta_W = g' cos\theta_W = \frac{gg'}{\sqrt{g^2 + g'^2}} \tag{3.13}$$

and

$$Q = \tau^3 + Y \tag{3.14}$$

for a fact we know that equations 3.11,3.12 and 3.14 are representation independent, hence they can applied in any representation.

finally the results imply that $\tilde{L}^a \nu \neq 0$ for a 1,2,3, while $\tilde{L}^4 \nu = 0$. That is$\tilde{L}^4$ unbrocken generator which we know as ieQ. indeed, $SU(2) \times U(1)_Y$ is spontaneously brocken down to $U(1)_{EM}$.

3.2 Baryon number violation and C,P and CP violation in the Standard Model

3.2.1 The B+L Anomaly

Many physics models other than the standard model predict novel sources of B violation, as we know. However, the baryon number is also violated in SM itself . The non-perturbative transtions between vacuum values of $SU(2)_L$ gauge theory change the combinations of charge (B+L), which can be described as subtle. Quantum tunnelling can prepare the path for these transitions at zero temperature at an unobservably slow pace, but it can also happen in the hot universe at $T > 100GeV$ through thermal fluctuations. In any case, rapid (B+L) violation occurs at high temperatures and

conservation occurs at low temperatures.

The violation of (B+L) is crucial in the mechanism of baryogenesis ($T > 100GeV$). First and foremost, we must note that the classical standard model Lagrangian is unaffected by the $U(1)B$ and $U(1)L$ transformations. These two are referred to as accidental symmetries, and they are both good up to the classical level. Session 3.1 demonstrates the dissection. The path integral measure we use to define the theory must also be invariant for these changes to have full quantum symmetry, but it isn't. As a result, a one loop calculation gives us

$$\partial_\mu j^\mu_B = \partial_\mu j^\mu_L = \frac{n_g}{32\pi^2}\left(g^2 W^a_{\mu\nu}\tilde{W}^{a\mu\nu} - g'^2 B_{\mu\nu}\tilde{B}^{\mu\nu}\right) \tag{3.15}$$

where n_g=3 which is the number of generations and $W^a_{\mu\nu}$, $B_{\mu\nu}$ are the respective field strength tensors of $SU(2)_L$ and $U(1)_Y$ and dual tensor $\tilde{W}^{a\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\alpha\beta}W^a_{\mu\nu}$.

From the classical symmetry we can get the expressions for Noether currents which is summed over the generations and it is give as

$$j^\mu_B = \frac{1}{3}\sum_i\left(\tilde{Q}^i_L\gamma^\mu Q^i_L + \tilde{u}^i_R\gamma^\mu u^i_R + \tilde{d}^i_R\gamma^\mu d^i_R\right) \tag{3.16}$$

$$j^\mu_L = \sum_i\left(\tilde{L}^i_L\gamma^\mu L^i_L + \tilde{e}^i_R\gamma^\mu e^i_R\right) \tag{3.17}$$

from equation 3.15 we can notice one that it can be written as total divergence:

$$\frac{g^2}{32\pi^2}W^a_{\mu\nu}\tilde{W}^{a\mu\nu} = \partial_\mu K^\mu \tag{3.18}$$

with

$$K^\mu_2 = \frac{g^2}{32\pi^2}\epsilon^{\mu\nu\alpha\beta}\left(W^a_{\nu\alpha}W^a_\beta - \frac{g}{3}\epsilon^{cab}W^c_\nu W^a_\alpha W^b_\beta\right) \tag{3.19}$$

where $\epsilon^{cab} = f^c ab$ is know as $SU(2)_L$ structure constant. similarly for $U(1)_Y$ the expression is $f^{cab} \rightarrow 0$.

The non-conservation of (B+L) is described by equation 3.15 , although it is unclear how it arises. Let's look at the vacuum structure of pure non-abelian gauge theories for a better understanding. It can be demonstrated that such theories contain a large number of independent classical vacua that can be labelled with integers $N_{CS} \in Z$.

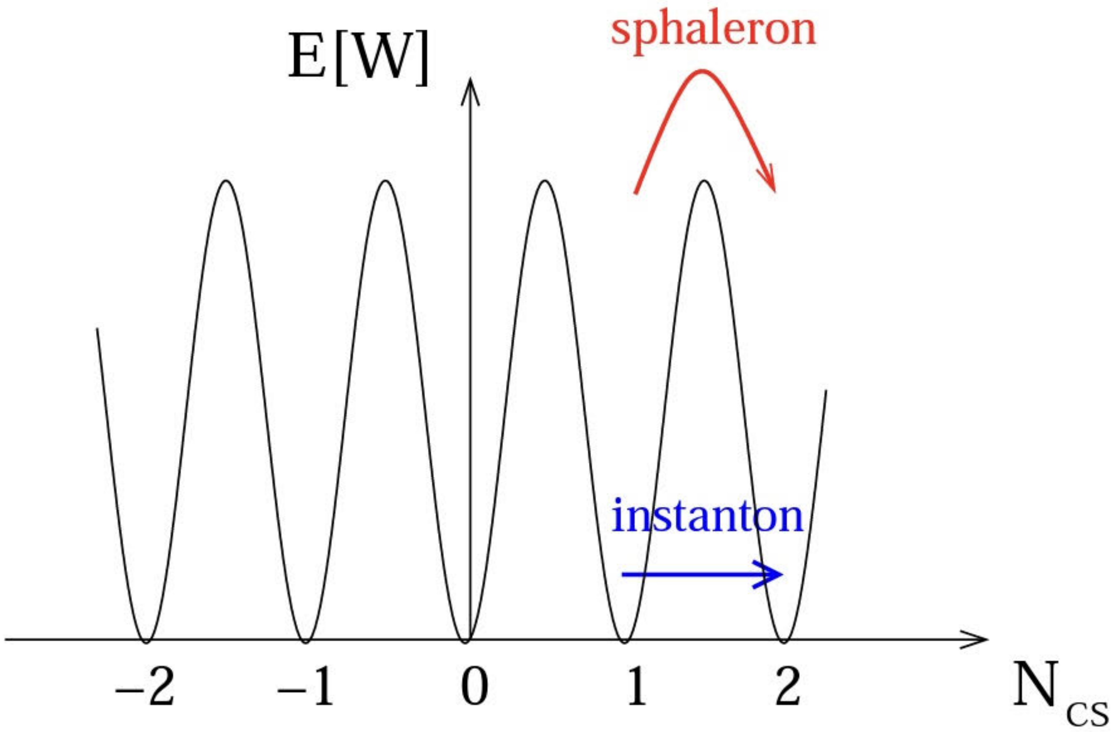

Figure 3.1: Vacuum structure of a non-Abelian gauge theory.

(refer the figure 1). These integers are also known as the Chern-Simons number of the mapping because they correspond to the topology of the mapping of a non-abelian group to the Euclidean spacetime boundary at infinity. Let us define the time-gauge dependent quantity $N_{CS}(t)$ as

$$N_{CS}(t) = \int d^3xK_2^0 \tag{3.20}$$

$$\rightarrow -\frac{g^3}{96\pi^2}\int d^3x\epsilon^{ijk}W_{\cdot i}^{a}W_{j}^{b}W_{k\cdot}^{c}.(W_0^a \rightarrow 0gauge, W_{\mu\nu}^a = 0) \tag{3.21}$$

When this classical vacuum configuration is evaluated with $W_0^a = 0$ (achieved through gauge choice) and $W_{\mu\nu}^a$ (such that it is a vacuum configuration), it becomes Time-Independent and coincides with the vacuum's N_{CS} value. Apart from having a non-trivial vacuum structure, it is obvious that there are solutions to the Euclidean space non-abelian field equations with finite actions that connect vacua with distinct N_{CS}

values. Because they have a finite extent in both space and time, these solutions are called Instantons.

Let us compute the change in the charge $B = \int d^3x j_B^0$ from $t = -\infty$ to $t = \infty$

$$\Delta B = \int_{-\infty}^{\infty} dt\partial_0 \int d^3x j_B^0 = \int_{-\infty}^{\infty} dt \int d^3x \left[\bar{\nabla}.\bar{j}_B + \frac{n_g}{32\pi^2}\left(g^2 W_{\mu\nu}^a \widetilde{W}^{a\mu\nu} - g'^2 B_{\mu\nu}\widetilde{B}^{\mu\nu}\right)\right] \tag{3.22}$$

the spatial gradient term above gives zero because it reduces to surface term involving the fermion field. Let us focus on the $SU(2)_L$ and evaluate it by using equation 3.18 with W_0^a=0,

$$\int d^4x \frac{g^2}{32\pi^2} W_{\mu\nu}^a \widetilde{W}^{a\mu\nu} = \int_{-\infty}^{\infty} dt \int d^3x (\partial_0 K_2^0 - \bar{\nabla}.\bar{K}_2) \tag{3.23}$$

$$= \int d^3x K_2^0|_{t\to\infty} - \int d^3x k_2^0|_{t\to-\infty} + 0 \tag{3.24}$$

$$= N_{CS}(t \to \infty) - N_{cs}(t \to -\infty) \tag{3.25}$$

The field strengths must vanish at spacetime infinity in order to identify the spacial integral of K^0 with N_{CS} and establish $K^i \to 0$ on the spatial boundary. The same reasoning prove that the hypercharge gauge field's corresponding expression vanishes.

By combining equation 3.22 with equation 3.22, we get

$$\Delta B = n_g \Delta N_{CS} \tag{3.26}$$

As a result, we may deduce that changes in the vacuum state of $SU(2)_L$ non-abelian gauge theory correlate to baryon number violation. It can also be demonstrated that the euclidean action for such a tunneling(occurring through quantum tunnelling transitions termed instantons) event is $S_E \geq 8\pi^2|\Delta N_{CS}|/g^2$, implying that the tunnelling rate is proportional to

$$\Gamma \propto e^{S_E} \leq e^{-8\pi^2/g^2} \simeq 10^-160 \tag{3.27}$$

this dimensionless factor makes the instanton transitions much too slow to be observed and consistent with the stability of proton.

At a fixed temperature, the situation is considerably different. Thermal fluctuations

can drive the system over top of the barriers between vacua rather than tunnelling through them. The sphaleron is a configuration that includes both the vector fields W_μ^a and the Higgs field H. Lattice studies, which offer the best estimate of the rate, are the most accurate.

$$\Gamma_{sp} = (18 \pm 3)\alpha_W^5 T^4; \qquad \alpha_w = \frac{g^2}{4\pi}$$

ie, Since B-L is conserved, the change in B due to gauge transformations accounts for

$$\Delta B = \frac{1}{2}\Delta(B + L) = n_g \Delta N_{CS}$$

3.2.2 C,P and CP on the SM

To study the effect of Charge Conjugation symmetry (C) and the product of charge conjugation and parity(CP), let us begin with a complex scalar field ϕ and the Lagrangian associated with it given as

$$\mathcal{L} = |\partial\phi|^2 - m^2|\phi|^2 \tag{3.28}$$

we know that it has a Global $U(1) \cong SO(2)$ global symmetry, with the corresponding Noether current

$$j_\mu = i\phi^\dagger \overleftrightarrow{\partial_\mu} \phi \tag{3.29}$$

if we expand in the free-field limit, we get

$$\phi(x) = \int \widetilde{dk}\Big[a(\bar{k})e^{-ik.x} + b^\dagger(\bar{k})e^{ik.x}\Big] \tag{3.30}$$

$$Q = \int d^3x j^0 = \int \widetilde{dk}\Big[a^\dagger(\bar{k})a(\bar{k}) - b^\dagger(\bar{k})b(\bar{k})\Big] \tag{3.31}$$

where $\widetilde{dk} = d^3x/2E_k(2\pi)^3$. Let us identify $a^\dagger(\bar{k})$ as the creation operator for particle with positive charge and $b^\dagger(\bar{k})$ for a particle with negative charge.we call the b-species the antiparticle of the a-species because the theory has global $U(1)$ symmetry and it forces the masses of these particle to be equal and their charges to be opposite. The natural definition of charge conjugation(C) on this theory is the exchange of particles

and antiparticles in the sense

$$a(\bar{k}) \leftrightarrow b(\bar{k}) \tag{3.32}$$

similarly in terms of the field operators

$$\phi(x) \leftrightarrow \phi^{\dagger}(x) \tag{3.33}$$

this implies $Q \to -Q$ and it is clear that the Lagrangian is invarient under the transformation.It's also worth noting that in the free theory, the mode of expansion for a dirac fermion is given by the four-component chiral representation, which is written as,

$$\psi(x) = \int \widetilde{dk} \sum_{s} \left[a_s(\bar{k}) u(k,s) e^{-ik.x} + b_s^{\dagger}(\bar{k}) v(k,s) e^{ik.x} \right] \tag{3.34}$$

the standard free dirac theory has a built in U(1) symmetry with a current $j^{\mu} = \bar{\psi}\gamma^{\mu}\psi$ and charge

$$Q = \int \widetilde{dk} \sum_{s} \left[a_s^{\dagger}(\bar{k}) a_s(\bar{k}) - b_s^{\dagger}(\bar{k}) b_s(\bar{k}) \right] \tag{3.35}$$

where $a_s \dagger (\bar{k})$ is the creation operator for a fermion with momentum $\bar{k}$ and spin s, and $b_s^{\dagger}(\bar{k})$ is the creation operator for the antifermion with the same momentum and spin as the fermion but opposite charge. As a result, the natural definition of C is

$$a_s(\bar{k}) \leftrightarrow b_s(\bar{k}) \tag{3.36}$$

in short the effect in the chiral representation can be shown as[24,9]

$$\psi(x) \to (-i\gamma^2\gamma^0)\bar{\psi}^t \tag{3.37}$$

$$\bar{\psi}(x) \to \psi^t(-i\gamma^2\gamma^0) \tag{3.38}$$

based on these results, define $\Gamma_C = -i\gamma^2\gamma^0 = -\Gamma_C^t = -\Gamma_C^{\dagger}$. The action of C on fermion bilnears is then[9]

$$\bar{\psi}\Gamma_i\chi \to \bar{\chi}(\Gamma_C^{-}1\Gamma_i^t\Gamma_C)\psi = \eta_i\bar{\chi}\Gamma_i\psi \tag{3.39}$$

where $\eta_i = +1$ for $\gamma_i = 1, i\gamma^5, \gamma^{\mu}\gamma^5$ and $\eta_i = -1$ for $\gamma_i = \gamma^{\mu}, \sigma^{\mu\nu}$. The current j^{μ} is odd under C as expected and the free dirac action is also invarient.

The way the photon field interacts to EM current in QED can be used to determine

the action of C on vector bosons.

$$\mathcal{L}_{QED} = -eA_\mu \sum_i Q_i \bar{\psi}_i \gamma^\mu \psi_i \tag{3.40}$$

Because we know that the current is odd under C, we may write $j^\mu_{em} \to -j^\mu_{em}$ and the theory will be invariant under C, implying that the photon field is also odd.

In QCD, it is comparable with one key difference. The current of matter is provided by,

$$j^{\mu a}_{QCD} = \sum_i \bar{q}_i \gamma^\mu t^a q_i \to -\sum_i \bar{q}_i \gamma^\mu (t^a)^t q_i \tag{3.41}$$

where t^a is a generator of the fundamental representation of $SU(3)_c$.It is permissible for the to select generators that are either symmetric or antisymmetric given the values of 'a'. Consequently, this indicates that

$$j^{\mu a}_{QCD} \to -\eta(a) j^{\mu a}_{QCD} \tag{3.42}$$

value of $\eta(a)$ is +1/-1 depending if t^a is symmetric/antisymmetric repectively. one can show that this is consistent with the gluon self interactions of QCD provided we assign

$$G^a_\mu \to -\eta(a) G^a_\mu \tag{3.43}$$

QCD field strength $G^a_{\mu\nu}$ is also odd up to the $\eta(a)$ factor when this is taken into account. Both QCD and QED are invariant under C, as we can see from these formulations. We must consider chiral fermions coupling to vector bosons in order to apply C to the whole Standard model. Let ψ be

$$\psi = \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix} \tag{3.44}$$

by using the action of C on the defined chiral components, we get

$$\psi_L \to i\sigma^2 \psi^*_R \tag{3.45}$$

$$\psi_R \to -i\sigma^2 \psi^*_L \tag{3.46}$$

Along with this, we define C's action on electroweak vector bosons in the same way that

we define it for gluons and photons. Because the chiral ψ_L and ψ_R components of the SM fermions are derived from distinct representations of the underlying $SU(2)_L \times U(1)_Y$ gauge group, the SM breaks C.

Now Parity(P), let us recall that it acts on spacetime as $t \to t$ and $\bar{x} \to -\bar{x}$. then for a scalar field this becomes

$$\phi(t, \bar{x}) \to \phi(t, -\bar{x}) \tag{3.47}$$

and for dirac fermion it becomes

$$\psi(t, \bar{x}) \to \gamma^0 \psi(t, -\bar{x}) \tag{3.48}$$

As a result, under P, the spatial components of the QCD and QED matter currents are odd, implying that parity is preserved in these theories if

$$A^a_\mu \to -\eta(\mu) A^a_\mu(t, -\bar{x}) \tag{3.49}$$

with $\eta(\mu = 0) = -1$ and $\eta(\mu = i) = 1$. Then for the chiral fermions , the transformation law implies as

$$\psi_L(t, \bar{x}) \to \psi_R(t, -\bar{x}) \tag{3.50}$$

$$\psi_R(t, \bar{x}) \to \psi_L(t, -\bar{x}) \tag{3.51}$$

Similarly to C, the chiral fermion representations under the electroweak group violate P in the SM. While the presence of chiral fermions in SM breaks C and P, the combination of CP does not. When these operations are combined, we get a CP transformation on the Dirac fermion ψ as

$$\psi(t, \bar{x}) \to -\Gamma_C \psi^*(t, \bar{x}) \tag{3.52}$$

it can be shown in terms of chiral components as given below,

$$\psi_L(t, \bar{x}) \to -i\sigma^2 \psi^*_L(t, \bar{x}) \tag{3.53}$$

$$\psi_R(t, \bar{x}) \to i\sigma^2 \psi^*_R(t, \bar{x}) \tag{3.54}$$

Each chiral component consistently changes into itself. The CP transformation essentially replaces each operator with its complex conjugate while leaving the numerical

coefficients unchanged. Let us take an example,

$$\bar{e}_R i\gamma^\mu(\partial_\mu - ig' B_\mu)e_R \to \bar{e}_R i\gamma^\mu(\partial_\mu - ig' B_\mu)e_R \tag{3.55}$$

$$(y_U)_{ij}\bar{Q}_{Li}Hu_{Rj} + (y_U^\dagger)_{ji}\bar{u}_{Rj}H^\dagger Q_{Li} \to (y_U^t)_{ji}\bar{u}_{Rj}H^\dagger Q_{Li} + (y_U^*)_{ij}\bar{Q}_{Li}Hu_{Rj} \tag{3.56}$$

If the Yukawa matrix is real, the operator in equation 3.56 is invariant, while the operator in equation 3.55 is clearly invariant. Equation 3.56 yields a result that is more generic than the SM: The presence of complicated couplings in the activity is usually connected with CP violation. While the presence of complex couplings in an action can lead to CP violation, it does not always do so. Let's have a look at how this works in the SM. The fermion fields are frequently modified after electroweak symmetry breaking by making field redefinitions in flavour space to generate diagonal mass matrices with positive mass eigenvalues. The Cabbibo-Kobayashi-Maskawa (CKM) matrix is the only residue of these transformations in the theory's interactions. The flavour transformations, to be precise, take the form of

$$u_L \to V_{u_L}U_L, u_R \to V_{u_R}U_R, d_L \to V_{d_L}d_L, d_R \to V_{d_R}d_R \tag{3.57}$$

which give the CKM matrix

$$V_{CKM} = V_{u_L}^\dagger V_{d_L} \tag{3.58}$$

This is a 3×3 unitary matrix with three rotation angles and six phases that can be parametrized. CP violation could be caused by the phases in the CKM matrix. However, not all of these phases are physical because we can make more field redefinitions that affect the CKM matrix but not the positive diagonal fermion mass matrices. These have the form

$$V_{u_L} = diag(e^{i\alpha_1}, e^{i\alpha_2}, e^{i\alpha_3}) = V_{u_R} \tag{3.59}$$

$$V_{d_L} = diag(e^{i\beta_1}, e^{i\beta_2}, e^{i\beta_3}) = V_{d_R} \tag{3.60}$$

There are six stages here, and they can be utilised to eliminate five of the CKM matrix's six independent phases. In the CKM, this results in a single irreducible phase, which causes physical CP violation in the SM. This counting also suggests that a CKM matrix with $(ng-1)(ng-2)/2$ physical CP violating phases would exist for a version of the SM with ng generations. CKM phase CP violation in Kaon and B-meson mixing and decay has been measured in experiments. par In the SM, there is another potential cause of

CP violation that poses a severe puzzle. It corresponds to the permitted operator.

$$L_{SM} \supset \frac{g_3^2}{32\pi^2}\Theta G^a_{\mu\nu}\widetilde{G}^{a\mu\nu} \tag{3.61}$$

This operator is inconsistent with P (and T, but not C), and so with CP. Its origin may be traced back to $SU(3)_c$'s nontrivial vacuum structure, which reflects a quantum superposition of the classical N_CS vaccua . A T-violating permanent electric dipole moment (EDM) of the neutron is a potentially observable result of a non-zero Θ. Current limits on the neutron EDM suggests

$$\Theta \leq 10^{-10} \tag{3.62}$$

The puzzle of why this dimensionless parameter (with natural range $\Theta \in [0,\, 2\pi]$) is so small is known as the strong CP problem. It's also worth noting that Θterm can be written for $SU(2)_L$. The $SU(2)_L$ term, unlike $SU(3)_c$ (with large quarks), is not observable since it can be removed by applying a (B + L) transformation to the SM action without changing anything else.

While the SM clearly breaks C, P, and CP, the combination CPT is respected. For any acceptable Lorentz quantum field theory, this is a highly universal need.

CHAPTER 4

BARYOGENESIS MECHANISMS

Contents

4.1 Baryogenesis Mechanisms

4.1.1 Electroweak Baryogenesis

Electroweak Baryogenesis (EWBG) is a set of methods for producing baryons at the electroweak phase transition in the early universe, when the Higgs field changes from $\langle H^\dagger H\rangle = 0$ to $\langle H^\dagger H\rangle = v^2 \simeq (174GeV)^2$. If the transition is strongly first order, it occurs through the nucleation of broken phase bubbles inside a symmetric plasma environment. These bubbles expand, collide, and combine until the cosmos is completely filled. In EWBG, baryon formation takes place along the bubble walls, which cause a significant divergence from thermodynamic equilibrium. Net chiral asymmetries (e.g.

more left-handed quarks than right-handed quarks) can be generated by particle scattering off the bubble walls with C and CP violations, which bias the sphaleron transitions outside the bubbles to produce more baryons than antibaryons. The baryons are swiftly carried into the inside of the expanding bubbles, where they are essentially stable.

4.1.2 GUT Baryogenesis

The GUT baryogenesis models, which have also been attempted in the context of SUSY GUTS , are the oldest baryogenesis models. You require baryon number violation if you start with a Universe that is baryon number symmetric and wish to end up with baryon asymmetry in the Universe. Heavy gauge and Higgs bosons interact in GUTS, and their interactions violate baryon number. Consider a heavy boson X with two decay channels $X \rightarrow q\bar{l}$,$q\bar{q}$, and decay rates r_1 and r_2, respectively. Let $\bar{X}$ be its antiparticle, with comparable decay channels and decay rates of $r_1^{'}$ and $r_2^{'}$ respectively. Let q and l be quarks and leptons, respectively, with baryon numbers of 1/3 and 0. If one decay mode is used to assign a baryon number to X, the other decay mode will violate the baryon number. As a result, one has a violation of baryon number. The decay rate r_1 of X to quarks must not be identical to the decay rate $r_1^{'}$ of $\bar{X}$ to antiquarks, which implies CP violation. Finally, to establish a baryon asymmetry, B violating interactions must be out of thermal equilibrium, because in equilibrium, the particle interactions will push the baryon number to 0. The three prerequisites are known as Sakharov's conditions: B violation, CP violation, and the out-of-equilibrium requirement. GUT baryogenesis theories have a number of flaws.

1)The mass of the X particle is usually around 10^{16} GeV. However, in an inflationary Universe, the mass of the inflaton is on the order of 10^{13} GeV (based on CMBR anisotropy measurements), therefore creating X particles during reheating via direct inflaton decays or in the thermal bath is problematic

2)When the mass ofX is less than 10^{13} GeV, proton decay becomes a concern, which is worsened for supersymmetric GUTs.

3)There are a number of massive scalars in supergravity and superstring theories that decay after the GUT phase transition, releasing a tremendous quantity of energy into

the Universe. The baryon-to-entropy ratio is diluted as a result.

4.1.3 Affleck-Dine Baryogenesis

The heaviness of the Higgs boson is the major impediment to the Standard Model's ability to explain the universe's baryon asymmetry. The required new physics may exist at either a very high or very low scale. Inflation most likely occurred in the early cosmos, according to a growing amount of evidence. As a result, baryogenesis must have occurred during or after reheating. The reheat temperature should not surpass 109 GeV to prevent overproducing weakly interacting light particles, such as gravitinos and other new states anticipated in theory. GUT baryogenesis is hampered by this. This also restricts leptogenesis options. On the other hand, Affleck-Dine baryogenesis is compatible with the low energy and temperature scales needed for inflation. Affleck-Dine baryogenesis, the last low-reheat SUSY scenario, may naturally recreate the universe's observable baryon asymmetry. Many particle physics models result in substantial entropy generation at late periods (Cohen, Kaplan and Nelson, 1993). This reduces the number of baryons in the system. Coherent generation can be quite efficient, and in many models, it is exactly this late dilution that results in today's low baryon density.

4.1.4 Leptogenesis

Electroweak sphaleron transitions reprocess an imbalance in lepton number L into a non-zero B asymmetry during leptogenesis. The decay of an extremely heavy SM-singlet neutrino is a very interesting hypothesis for generating the L asymmetry. Many theories for the masses and mixings of the neutrinos we detect contain such neutrinos, and they can cause explicit L violation. They can induce an initial lepton imbalance Li if their decays also include C and CP violation and are out of equilibrium. Sphaleron transitions then convert some of this lepton charge to baryon charge, resulting in the observed baryon asymmetry. Consider the case when we have a lepton imbalance of 10 as a result of heavy neutrino decays. If we rewrite this in terms of B+L and B-L, we get B+L=10 and B-L=-10, assuming that there is no initial net baryon number in the Universe.If we have B+L violating processes in thermal equilibrium, the B+L

will be reduced to 0 but the B-L will remain constant. As a result, B=-5 and L=5 are obtained. As a result, the Universe's lepton imbalance has been changed to a baryon asymmetry. Detailed discussion of leptogenesis along with the neutrino sector is given in the next section.

4.2 Inflation and Affleck-Dine Scenario

An angular motion of a complex scalar field phi charged under an approximation U(1) symmetry produces asymmetry. The dynamics of may result in an asymmetry number density linked to its U(1) charge.

$$n_\phi = 2Q \,\mathrm{Im}\left[\phi^\dagger \dot{\phi}\right],$$

which satisfies the equation of motion,

$$\dot{n}_\phi + 3Hn_\phi = \mathrm{Im}\left[\phi \frac{\partial V}{\partial \phi}\right] \tag{4.1}$$

A baryon number asymmetry can be created prior to the Electroweak Phase Transition if the U (1) symmetry is made up of the global U (1)B or U(1)L symmetries (EWPT). We'll look at a triplet Higgs that's doubly charged under the lepton number symmetry U(1)L in the next section. There are interaction terms between the triplet and the SM Higgs that clearly contradict lepton number, implying that Leptogenesis is conceivable.

Let us consider the lagrangian containing Higgs Doublet H and triplet Higgs Δ,

$$\begin{aligned}\frac{\mathcal{L}}{\sqrt{-g}} =& -\frac{1}{2}M_P^2 R - f(H,\Delta)R - g^{\mu\nu}\left(D_\mu H\right)^\dagger \left(D_\nu H\right) \\ & - g^{\mu\nu}\left(D_\mu \Delta\right)^\dagger \left(D_\nu \Delta\right) - V(H,\Delta) + \mathcal{L}_{\text{Yukawa}}\,,\end{aligned}$$

where $f(H,\Delta) = \xi_H H^\dagger H + \xi_\Delta \Delta^\dagger \Delta + \ldots$, M_p is the re- duced Planck mass, and R is the Ricci scalar. The Higgs' H and Δ are parameterized by,

$$H = \begin{pmatrix} h^+ \\ h \end{pmatrix}, \Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix},$$

where h and Δ^0 denotes the neutral components of H and Δ respectively. The $\mathcal{L}_{\text{Yukawa}}$ term contains not only the Yukawa interactions of the SM fermions, but also a new interaction between the left-handed leptons and the triplet Higgs Δ,

$$\mathcal{L}_{Yukawa} = \mathcal{L}_{\text{Yukawa}}^{\text{SM}} - \frac{1}{2} y_{ij} \bar{L}_i^c \Delta L_j + \text{ h.c.} \tag{4.2}$$

As the Δ^0 obtains a non-zero vacuum expectation value (VEV),the neutrino mass matrix will be generated through this interaction term. The triplet Higgs is given a lepton charge of $Q_L = -2$ as a result of this interaction, raising the prospect that it may play a role in the origin of baryon asymmetry. The potential for the Higgs' $V(H, \Delta)$ is given by,

$$\begin{aligned} V(H,\Delta) =& - m_H^2 H^\dagger H + \lambda_H \left(H^\dagger H\right)^2 + m_\Delta^2 \operatorname{Tr}\left(\Delta^\dagger \Delta\right) \\ &+ \lambda_3 \operatorname{Tr}\left(\Delta^\dagger \Delta\right)^2 + \lambda_1 \left(H^\dagger H\right) \operatorname{Tr}\left(\Delta^\dagger \Delta\right) \\ &+ \lambda_2 \left(\operatorname{Tr}\left(\Delta^\dagger \Delta\right)\right)^2 + \lambda_4 H^\dagger \Delta \Delta^\dagger H \\ &+ \left[\mu \left(H^T i\sigma^2 \Delta^\dagger H\right) + \frac{\lambda_5}{M_P}\left(H^T i\sigma^2 \Delta^\dagger H\right)\left(H^\dagger H\right)\right. \\ &\left. + \frac{\lambda_5'}{M_P}\left(H^T i\sigma^2 \Delta^\dagger H\right)\left(\Delta^\dagger \Delta\right) + h.c.\right] + \ldots \end{aligned}$$

the terms appeared in the square brackets violate lepton number and the dimension five operators are suppressed by M_P. Focusing on the neutral components of Δ and H(non-trivial VEVs),therefore

$$\begin{aligned} V\left(h,\Delta^0\right) =& -m_H^2 |h|^2 + m_\Delta^2 \left|\Delta^0\right|^2 + \lambda_H |h|^4 + \lambda_\Delta \left|\Delta^0\right|^4 \\ &+ \lambda_{H\Delta} |h|^2 \left|\Delta^0\right|^2 \\ &- \left(\mu h^2 \Delta^{0^*} + \frac{\lambda_5}{M_P}|h|^2 h^2 \Delta^{0^*} + \frac{\lambda_5'}{M_P}\left|\Delta^0\right|^2 h^2 \Delta^{0^*} + \text{ h.c. }\right) + \ldots \end{aligned}$$

where $\lambda_\Delta = \lambda_2 + \lambda_3$, and $\lambda_{H\Delta} = \lambda_1 + \lambda_4$ All of the parameters must satisfy the stability condition.non-vanishing Δ^0 VEV can be approximated in the limit $M_\Delta \gg v_{\mathrm{EW}}$ as,

$$\left\langle\Delta^0\right\rangle \simeq \frac{\mu v_{\mathrm{EW}}^2}{2m_\Delta^2} \tag{4.3}$$

where the SM Higgs VEV is $v_{\mathrm{EW}} = 246\mathrm{GeV}$. Note that the Δ^0VEV is bounded by $\mathcal{O}(1)\mathrm{GeV} > \langle\Delta^0\rangle$ 0.05 eV.By ensuring y_ν is perturbative up to M_p. in order to generate the observed neutrino masses The upper bound on the Δ^0 VEV is derived from T-parameter constraints determined by precision measurements.

Furthermore, the inflationary scenario will be caused by both Higgs through their non-minimal gravitational couplings. For high field values, these couplings act to flatten the scalar potential. This type of inflationary process has been used to produce

a Starobinsky-like inflationary era in standard Higgs inflation. The following non-minimal coupling is considered:

$$f(H,\Delta) = \xi_H |h|^2 + \xi_\Delta \left|\Delta^0\right|^2 = \frac{1}{2}\xi_H \rho_H^2 + \frac{1}{2}\xi_\Delta \rho_\Delta^2 \tag{4.4}$$

where we have utilized the polar coordinate parametrization $h \equiv \frac{1}{\sqrt{2}}\rho_H e^{i\eta}, \Delta^0 \equiv \frac{1}{\sqrt{2}}\rho_\Delta e^{i\theta}$. Here the inflaton can then be defined as φ which is related to ρ_H and ρ_Δ through,

$$\begin{aligned}
&\rho_H = \varphi \sin\alpha, \rho_\Delta = \varphi\cos\alpha \\
&\xi \equiv \xi_H \sin^2\alpha + \xi_\Delta \cos^2\alpha
\end{aligned}$$

hence the following Lagrangian,

$$\begin{aligned}
\frac{\mathcal{L}}{\sqrt{-g}} =& -\frac{1}{2}M_P^2 R - \frac{1}{2}\xi\varphi^2 R - \frac{1}{2}g^{\mu\nu}\partial_\mu\varphi\partial_\nu\varphi \\
& - \frac{1}{2}\varphi^2\cos^2\alpha g^{\mu\nu}\partial_\mu\theta\partial_\nu\theta - V(\varphi,\theta)
\end{aligned}$$

where

$$V(\varphi,\theta) = \frac{1}{2}m^2\varphi^2 + \frac{\lambda}{4}\varphi^4 + 2\varphi^3\left(\tilde{\mu} + \frac{\tilde{\lambda}_5}{M_P}\varphi^2\right)\cos\theta \tag{4.5}$$

and

$$\begin{aligned}
&m^2 = \left(m_\Delta^2\cos^2\alpha - m_H^2\sin^2\alpha\right) \\
&\lambda = \lambda_H \sin^4\alpha + \lambda_{H\Delta}\sin^2\alpha\cos^2\alpha + \lambda_\Delta\cos^4\alpha \\
&\tilde{\mu} = -\tfrac{1}{2\sqrt{2}}\mu\sin^2\alpha\cos\alpha \\
&\tilde{\lambda}_5 = -\tfrac{1}{4\sqrt{2}}\left(\lambda_5\sin^4\alpha\cos\alpha + \lambda_5'\sin^2\alpha\cos^3\alpha\right)
\end{aligned}$$

We regard lepton asymmetry to be a dynamical field since its value is determined by the motion of θ. In the model, m denotes that during the inflationary era, the quartic potential term is dominant. As a result, we expect the dim-5 operator to provide the majority of the U(1) breaking during inflation. By using the transformation given below we can translate into einstein frame

$$\tilde{g}_{\mu\nu} = \Omega^2 g_{\mu\nu}, \quad \Omega^2 = 1 + \xi\varphi^2/M_P^2$$

now we define the field to get a canonical kinetic term χ in terms of ϕ,

$$\begin{aligned}\chi(\varphi) =& 1/\sqrt{\xi}\left(\sqrt{1+6\xi}\sinh^{-1}\left(\sqrt{\xi+6\xi^2}\varphi\right)\right.\\ &\left.-\sqrt{6\xi}\sinh^{-1}\left(\sqrt{6\xi^2}\varphi/\sqrt{1+\xi\varphi^2}\right)\right.\end{aligned}$$

Subsequently, obtaining the final Einstein frame Lagrangian,

$$\begin{aligned}\tfrac{\mathcal{L}}{\sqrt{-g}} &= -\tfrac{M_P^2}{2}R - \tfrac{1}{2}g^{\mu\nu}\partial_\mu\chi\partial_\nu\chi\\ &-\tfrac{1}{2}f(\chi)g^{\mu\nu}\partial_\mu\theta\partial_\nu\theta - U(\chi,\theta)\end{aligned}$$

where

$$f(\chi) \equiv \frac{\varphi(\chi)^2\cos^2\alpha}{\Omega^2(\chi)}, \text{ and } U(\chi,\theta) \equiv \frac{V(\varphi(\chi),\theta)}{\Omega^4(\chi)}$$

hence from the Lagrangian we get,

$$\ddot{\chi} - \frac{1}{2}f'(\chi)\dot{\theta}^2 + 3H\dot{\chi} + U_{,\chi} = 0 \tag{4.6}$$

$$\ddot{\theta} + \frac{f'(\chi)}{f(\chi)}\dot{\theta}\dot{\chi} + 3H\dot{\theta} + \frac{1}{f(\chi)}U_{,\theta} = 0 \tag{4.7}$$

Let us assume that the model parameters are such that the dynamics of have a negligible impact on the inflationary trajectory, which will be justified further down. The resulting inflationary observables are identical to those of the Starobinsky model, and they fit well with existing observational constraints.

Lepton Asymmetry Considering the inflation dynamics, lepton number density

$$n_L = Q_L\varphi^2\dot{\theta}\cos^2\alpha$$

suppose that the initial $\dot{\theta}_0$is zero, and that the initial$\theta = \theta_0$ is non-zero. It will become clear that the sign of the resulting asymmetry is determined by the value of θ_0 chosen. The scale μ will be selected to be small enough that it does not contribute to lepton number violation during inflation and reheating, with the dim-5 term being the primary source.

$$\ddot{\theta} + \left(3H + \frac{\dot{f}(\chi)}{f(\chi)}\right)\dot{\theta} + \frac{2\tilde{\lambda}_5}{M_p}\frac{\phi(\chi)^5}{f(\chi)\Omega^4(\chi)}\sin\theta = 0 \tag{4.8}$$

Under these initial conditions, θ will be dominated by the potential and acceleration term. Due to the balance between Hubble friction and the slope of the potential, once a non-zero $\dot{\theta}_0$ is created, a terminal velocity is rapidly attained, driving the acceleration to zero. The resultant inflationary $\dot{\theta}$ is defined by the following equilibrium relation,

$$3H_{\text{inf}}\dot{\theta}_{\text{inf}} \approx \frac{2\tilde{\lambda}_5 M_p^2 \sin\theta_0}{\xi^{\frac{3}{2}} \cos^2\alpha} \sqrt{e^{\sqrt{\frac{2}{3}}\frac{\chi}{M_P}} - 1}$$

The oscillatory reheating period begins after inflation ends ($\chi \sim M_p$), and non-zero acceleration is generated by the χ oscillations.

$$\ddot{\theta} + \frac{f'(\chi)}{f(\chi)}\dot{\theta} \approx 0 \tag{4.9}$$

therefore

$$\dot{\theta}_{\text{reh}} \approx \frac{2\tilde{\lambda}_5 \sin\theta_0}{3M_p H \cos^2\alpha} \frac{\varphi_{\max}^3}{1 + \frac{\xi\varphi_{\max}^2}{M_p^2}} \tag{4.10}$$

so we can write the lepton number density

$$n_L^{\text{reh}} \approx Q_L \frac{2\tilde{\lambda}_5 \sin\theta_0}{3M_p H_{\text{reh}}} \varphi_{\text{reh}}^5$$

From the total energy density, $\frac{\lambda}{4}\varphi_{\text{reh}}^4 \simeq 3M_p^2 H_{\text{reh}}^2$. ,the approximate value of φ_{reh} can be found.

$$n_L^{\text{reh}} \approx Q_L \frac{15\tilde{\lambda}_5 \sin\theta_0}{\lambda^{\frac{5}{4}}} \left(M_p H_{\text{reh}}\right)^{\frac{3}{2}} \tag{4.11}$$

This relation agrees with numerical simulations within an order of magnitude.

Baryon Asymmetry: In this scenario, the reheating temperature is predicted to be high enough for this to happen due to SM Higgs couplings, $T_{reh} \sim 10^{13} - 10^{14}$ GeV, as in Higgs inflation. To get the anticipated baryon asymmetry parameter, assuming no substantial entropy generation after reheating $s = \frac{2\pi}{45} g_* T_{reh}^3$.

$$\frac{\eta_B}{\eta_B^{\text{obs}}} \simeq 1.7 \times 10^{10} \frac{\left|\tilde{\lambda}_5 \sin\theta_0\right|}{\lambda^{\frac{5}{4}}} \simeq 7 \times 10^{21} \frac{\left|\bar{\lambda}_5 \sin\theta_0\right|}{\xi^{\frac{5}{2}}} \tag{4.12}$$

where the observed baryon asymmetry is given as $\eta_B^{\text{obs}} \simeq 8.5 \times 10^{-11}$, and $H_{\text{reh}}^2 \simeq \frac{\pi^2}{90} g_* \frac{T_{-mb}^4}{M_p^2}$. therefore we can say the final asymmetry has a parameter dependence and

is independent of the reheating temperature. For the reference values $\xi \sim 300, \lambda \sim 4.5 \cdot 10^{-5}$ $\theta_0 \sim 0.05$, and $\tilde{\lambda}_5 \sim 5 \cdot 10^{-15}$, we can obtain the observed baryon asymmetry parameter ratio, $\eta_B \sim \eta_B^{\text{obs}}$. Given that the $U(1)_L$ breaking term couplings are both very small, it is obvious to consider that they originate from a sphaleron field which carries $U(1)_L$ charge.

To prevent the dim-5 term from dominating inflationary dynamics

$$\tilde{\lambda}_5 \ll 6.25 \cdot 10^{-11} \xi^{5/2} e^{-\frac{\chi_o}{\sqrt{6}M_P}}$$

Considering all the constraints we get an upper bound on the baryon asymmetry,

$$\frac{\eta_B}{\eta_B^{\text{obs}}} \ll 5 \cdot 10^{10} \sin\theta_0$$

where $N = 50$ is the minimum value, with bigger e-foldings reducing this limit. It is required for, to guarantee that the mu term is insignificant during reheating.

$$\frac{\tilde{\mu}}{M_p} \ll \frac{\tilde{\lambda}_5}{\xi^2} \ll 6.25 \cdot 10^{-11} \sqrt{\xi} e^{-\frac{x_0}{\sqrt{6}M_p}}$$

at the end of reheating,the triplet will be thermalized , and we must consider possible wash-out processes. Firstly, we require that the processes $LL \leftrightarrow HH$ are not effective,

$$\Gamma = n\langle\sigma v\rangle \approx \frac{y^2\mu^2}{m_\Delta} < H|_{T=m_\Delta},$$

Also the Higgs Triplet will genarate the neutrino mass ,given as

$$m_\nu \simeq y\frac{\mu v^2}{2m_\Delta^2} \tag{4.13}$$

where m_ν should be at least the order of the largest neutrino mass 0.05eV. Combining this with the above relation, we obtain the mass limit $m_\Delta < 10^{12}$GeV for $m_\nu = 0.05$eV To explain the neutrino masses in our model, the triplet Higgs vacuum expectation value is generally tiny, thus this constraint has no phenomenological repercussions.

Q-balls that are stable are rare to develop. However, because to the very high neutrino Yukawa couplings we require, even if they can be produced during reheating, they would decay considerably earlier than the EWPT.

4.3 Leptogenesis and Origin of Neutrino mass matrix

4.3.1 Three Neutrino Mixing

All of the evidence for neutrino oscillations so far has been based on 3-flavour neutrino mixing in vacuum. This is the simplest neutrino mixing system that can explain for the currently available data on solar, atmospheric, reactor, accelerator neutrinos oscillations. The left-handed fields of the flavour neutrinos ν_e, ν_μ,and ν_τ, and in the CC weak interaction Lagrangian equation for the weak charged lepton current are linear combinations of the Left handed components of the fields of three large neutrinos νj

$$\mathcal{L}_{CC} = \frac{-g}{\sqrt{2}} \sum_{l=e,\mu,\tau} \bar{l_L}(x)\gamma_\alpha \nu_{lL}(x) W^{\alpha\dagger}(x) + h.c \tag{4.14}$$

$$\nu_{lL}(x) = \sum_{j=1}^{3} U_{lj}\nu_{jL}(x) \tag{4.15}$$

where U is the unitary neutrino mixing matrix, which have 3 parameterised angles with 1 or 3 CP violation phases.

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c12s_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times diag(1, e^{i\frac{\alpha_{21}}{2}}, e^{i\frac{\alpha_{31}}{2}}) \tag{4.16}$$

where $c_{ij} = \cos\theta_{ij}$, $s_{ij} = \sin\theta_{ij}$ and the angles $\theta_{ij} = [0, \pi/2], \delta = [0, 2\pi]$ is the dirac CP violation and α_{21}, α_{31} are two marojona CP violation phases.In the case of heavy Dirac neutrinos, the neutrino mixing matrix U is comparable to the CKM quark mixing matrix in terms of the number of mixing angles and CPV phases. The appearance of two more physical CPV phases in U if ν_j are Majorana particles is due to the latter's unique characteristics.

Also the parameters for characterising the 3 neutrino mixing are given below:

1)Angles $\theta_{12}, \theta_{23}, \theta_{13}$ 2)Nature of massive neutrinos ν_j-1 Dirac (δ) or 1 Dirac + 2 Marojona ($\delta, \alpha_{31}, \alpha_{21}$), CPV phases. 3) Masses of neutrinos m1,m2,m3.

Parameter	best-fit	3σ
Δm^2_{21} $[10^{-5}$ eV $^2]$	7.37	6.93 – 7.97
$\|\Delta m^2\|$ $[10^{-3}$ eV $^2]$	2.50 (2.46)	2.37 – 2.63 (2.33 – 2.60)
$\sin^2\theta_{12}$	0.297	0.250 – 0.354
$\sin^2\theta_{23}$, $\Delta m^2 > 0$	0.437	0.379 – 0.616
$\sin^2\theta_{23}$, $\Delta m^2 < 0$	0.569	0.383 – 0.637
$\sin^2\theta_{13}$, $\Delta m^2 > 0$	0.0214	0.0185 – 0.0246
$\sin^2\theta_{13}$, $\Delta m^2 < 0$	0.0218	0.0186 – 0.0248
δ/π	1.35 (1.32)	(0.92 – 1.99)
		((0.83 – 1.99))

With help of neutrino mixing matrix elements one can define the angles as,

$c^2_{12} \equiv \cos^2\theta_{12} = \frac{|U_{e1}|^2}{1-|U_{e3}|^2}$, $\quad s^2_{12} \equiv \sin^2\theta_{12} = \frac{|U_{e2}|^2}{1-|U_{e3}|^2}$

$\mathrm{s}^2_{13} \equiv \sin^2\theta_{13} = |U_{e3}|^2$, $\quad s^2_{23} \equiv \sin^2\theta_{23} = \frac{|U_{\mu 3}|^2}{1-|U_{e3}|^2}$,

$\mathrm{c}^2_{23} \equiv \cos^2\theta_{23} = \frac{|U_{\tau 3}|^2}{1-|U_{e3}|^2}$

It is possible to measure the the values of mixing angles for different systems and they all give best fit values with range of 3σ. The $\theta_{12} \simeq 35^0$ angle, also known as the solar mixing angle, is best estimated through observations of neutrinos produced by nuclear reactions in the sun. The atmospheric mixing angle was initially determined using neutrinos from cosmic ray showers in the atmosphere, and it is frequently referred to as the atmospheric mixing angle$\theta_{23} \simeq 45^0$ Neutrinos generated in nuclear reactors have recently been measured at$\theta_{13} \simeq 14^0$. The mass differences of neutrinos may also be calculated via oscillation measurements, with

$$\triangle m^2_{12} \simeq 7.6 \times 10^{15} eV^2, \qquad |m^2_{23}| \simeq 2.4 \times 10^{-3} eV^2. \tag{4.17}$$

As neutrino masses and mixing need a new physics, with the help of yukawa couplings let us add three gauge singlet right- handed neutrinos since it is easy to generate,

N_{Ri}=(1,1,0), i=1,2,3

$$-\mathcal{L} \supset -\lambda_{Ai}\bar{N}_{Ri}H \cdot L_A + (h.c) \tag{4.18}$$

where $H = (H^+, H^0)^t$,$L^A = (\nu_{LA}, e_{LA})^t$, $A \cdot B = \epsilon_{ab}A^aB^b$ with $\epsilon_{12} = \pm 1$ for $SU(2)_L$ doublets. soon after the electroweak symmetry breaking H generates a mass matrix with entries,

$$(m_\nu)_{AB} = (\lambda\nu)_{AB} \tag{4.19}$$

As a result we get 3 heavy massive Dirac neutrinos and the mass matrix connecting the leptons through W boson.The Type-I neutrino seesaw is a common variant on the simple image above. We may also include (diagonal) Majorana masses for the NR because they are gauge singlets.

$$-\mathcal{L} \supset \frac{1}{2}M_i(\overline{N^c_{Ri}})N_{Ri} + (h.c) \tag{4.20}$$

$N^c_R = -i\gamma^2\gamma^0\overline{N}^t_R$, hence we can write the schematic form of mass matrix as

$$M_\nu = \begin{pmatrix} 0 & \lambda\nu \\ \lambda\nu & M_N \end{pmatrix} \tag{4.21}$$

for the condition $M_n >> y_N\nu$ the 6 Marojona fermions get the mass eigenvalues as

$$m_\nu \simeq \frac{(\lambda\nu)^2}{M_N}, \qquad M_N \tag{4.22}$$

The Standard Model neutrinos are used to identify the three light states, whereas the three heavy neutrinos are usually singlets and difficult to detect. For $M_N \sim 10^{13}$GeV, the SM-like neutrinos have masses in the sub-eV range. Examining the EFT obtained by integrating out the really large right-handed neutrinos is equally instructive. Hence the operator generated is given as

$$-\mathcal{L} \supset \sum_i \frac{\lambda_{Ai}\lambda_{Bi}}{2M_i}(\overline{L^c_A} \cdot H)(L_B \cdot H) + (h.c) \tag{4.23}$$

This is the smallest non-renormalizable operator that can be constructed only from SM fields. It creates neutrino masses on the order of $m_\nu \sim \lambda^2\nu^2/M_N$ after electroweak symmetry breakdown, as predicted by the neutrino seesaw.

4.4 cosmological Baryon and Lepton Production

To produce a lepton asymmetry in the early cosmos, the most prevalent model of leptogenesis depends on the decays of the heavy right-handed neutrinos in the Type-I seesaw model. Both C and CP violations are required to create such an imbalance. With decay asymmetries in hand, we go on to the cosmological mechanisms in the Type-I seesaw model that produce and remove lepton and baryon asymmetries. L is produced by the CP-violating decays of heavy right-handed neutrinos, which is subsequently partially transformed to B via sphaleron transitions. Inverse decay and scattering processes, on the other hand, tend to erase the L charge generated in the decays, necessitating a departure from thermodynamic equilibrium to avoid this. In this part, we look at how electroweak sphalerons redistribute B and L charges in complete equilibrium, as well as the cosmological time evolution of these and other charges away from equilibrium.

4.4.1 Redistribution of Sphalerons

Electroweak sphalerons transitions, as previously stated, violate (B + L). Remember that they are active in the SM at a higher effective rate than Hubble for $T \in [130, 10^{12}]$ GeV. This temperature range encompasses the majority of the temperature range important to leptogenesis, and sphaleron transitions are critical in converting the L charge generated into a B charge. Let us have a look at this process at equilibrium.

Number density asymmetry of a particle of species ϕ at a high temparature $T >> m_i$ ia give as

$$n_\phi - n_{\bar{\phi}} \simeq \left\{ g_\phi \mu_\phi T^2/6 \text{ for fermion; } g_\phi \mu_\phi T^2/3 \text{ for boson.} \right.$$

where g_ϕ is the number of internal degrees of freedom . As a result, we can track particle asymmetries by looking at their chemical potentials. The chemical potentials are now restricted by the set of rapid reactions that can occur, according to thermodynamic equilibrium. This may be used to calculate the final B and L charges after an instantaneous L charge injection.

The SM gauge interactions are in equilibrium with $\mu_Y = \mu_W = \mu_g = 0$ for the corresponding vector bosons at the temperatures crucial for leptogenesis. This means that

all of the gauge multiplet's components have the same chemical potential. As a result, we'll write

$$\mu_{Q_i} = \mu_{u_{Li}} = \mu_{d_{Li}}, \qquad \mu_{l_i} = \mu_{\nu L_i} = \mu_{e_{Li}} \tag{4.24}$$

with the help of Yukawa interactions we can write it as,

$$\mu_{Q_i} + \mu_H - \mu_{u_i} = \mu_{Q_i} - \mu_H - \mu_{d_i} = \mu_{L_i} - \mu_H - \mu_{e_i} = 0 \tag{4.25}$$

where $\mu_{u_i}, \mu_{d_i}, \mu_{e_i}$ are right handed fermions. For the sphalerons of $SU(2)_L$ and $SU(3)_c$ kind we have,

$$0 = \sum_i (3\mu_{Q_i} + \mu_{li}), \qquad 0 = \sum_i (2\mu_{Qi} - \mu_{ui} - \mu_{di}) \tag{4.26}$$

by defining baryon and lepton chemical potential we can relate the above relations to B and L.

$$n_B = \mu_B T^2/6, \qquad n_L = \mu_L T^2/6 \tag{4.27}$$

n_B and n_L are the charge density asymmetries. Hence we write

$$\mu_B = \sum_i (2\mu_{Qi} + \mu_{di} + \mu_{ui}), \qquad \mu_L = \sum_i (2\mu_{Li} + \mu_{ei}) \tag{4.28}$$

from all these definition and relations given, we can formulate the relationship between charges in standard model equilibrium as

$$\mu_B = c(\mu_B - \mu_L), \qquad \mu_L = (c-1)(\mu_B - \mu_L) \tag{4.29}$$

value of c can calculated as

$$c = \frac{(8n_g + 4)}{(22n_g + 13)} \rightarrow \frac{28}{79}, \qquad n_g = 3 \tag{4.30}$$

This conclusion, when combined with the chemical potential expressions, suggests that after equilibration, the final B density is proportional to the total (B-L) charge. Assume that the instantaneous decay of a heavy neutrino produces a lepton density asymmetry yield of $Y_L(t_i)$for leptogenesis. After sphaleron equilibration and other post-decay events, the ultimate baryon asymmetry yield today ($t = t_0$) will be equal to

$$Y_B(t_o) \simeq -c_s Y_L(t_i) \tag{4.31}$$

Successful baryogenesis necessitates a deviation from thermodynamic equilibrium, and this is reflected in the creation and decay of heavy neutrinos during leptogenesis. The evolution of the Ni and (B-L) number densities may be characterised by a series of semi-classical Boltzmann equations. After computing the yield $Y_{B-L}(z \to \infty)$, the baryon yield may be calculated using the result of Eq (5.16). The baryon asymmetry from leptogenesis is frequently expressed in the form when converted to a fraction relative to photons.

$$\eta \equiv \frac{n_B}{n_\gamma} \simeq -\mathcal{K}\epsilon_1 \tag{4.32}$$

where $\mathcal{K}$ is the efficiency factor and ϵ_1 corresonds to C and CP violation.

CHAPTER 5

BARYON AND LEPTON ASYMMETRY RELATIONS

Contents

5.0.1 Relation between Baryon and Lepton Asymmetry

Using the relationship between charges in equilibrium we can solve to get the relation between baryon and lepton asymmetry.(The notation $B \equiv \mu_B$ and $L \equiv \mu_L$ is often usued)

$$\mu_B = c_s(\mu_B - \mu_L) = B = c_s(B - L) = 0.35(B - L)$$

$$\mu_L = (c_s - 1)(\mu_B - \mu_L = L = (c_s - 1)(B - L) = -0.65(B - L)$$

hence;

$$\frac{\mu_B}{\mu_L} = \frac{0.35(B - L)}{-0.65(B - L)} = 0.538 = -0.54$$

where the chemical potential is defined as

$$n_B = \mu_B T^2/6, \qquad n_L = \mu_L T^2/6$$

using above parameters,we can write

$$\frac{\eta_1}{\eta_2} = \frac{n_B/n_\gamma}{n_L/n_\gamma} = \frac{(n_b - n_{\bar{b}})}{(n_l - n_{\bar{l}})} = \frac{\mu_B}{\mu_L} = \frac{B}{L} = \frac{c}{c-1} = -0.549 \approx -0.55$$

by the measurements of CMB , we know the value of η_1=6 $\times$ 10^{-10}. With the help above expression we can calculate the lepton asymmetry as

$$\frac{\eta_1}{\eta_2} = \frac{n_B}{n_L} = -0.55$$

$$\eta_2 = \frac{6 \times 10^{-10}}{-0.55} = -1.11 \times 10^{-9}$$

5.0.2 Comparison of present and initial yields of baryon and lepton Asymmetry

we know

$$\eta_1 = \frac{(n_b - n_{\bar{b}})}{n_\gamma} \qquad ;\eta_2 = \frac{(n_l - n_{\bar{l}})}{n_\gamma}$$

similarly the yield

$$Y_B = \frac{(n_b - n_{\bar{b}})}{s}; \qquad Y_L = \frac{(n_l - n_{\bar{l}})}{s}$$

if take the ratio of above expressions we get

$$\eta_1 = Y_B \frac{s}{n_\gamma}; \qquad \eta_2 = Y_L \frac{s}{n_\gamma}$$

if we consider the $Y_L(t_0)$ as the lepton asymmetry of today the we can write the ratio of η_1 and η_2 by considering

$$(Y_B(t_0) = -c_s Y_L(t_i))$$

$$\frac{\eta_1}{\eta_2} = \frac{Y_B}{Y_L} = -c\frac{Y_L(t_i)}{Y_L(t_0)}$$

using the value of ratio$\frac{\eta_1}{\eta_2}$ given in the previous section

$$\frac{\eta_1}{\eta_2} = \frac{c}{c-1} = -c\frac{Y_L(t_i)}{Y_L(t_0)}$$

hence we can write lepton asymmetry yield of present time to the initial time as

$$Y_L(t_0) = (1-c)Y_L(t_i)$$

$$Y_L(t_0) = 0.64 Y_L(t_i) \qquad ;Y_L(t_i) = \frac{1}{0.64}Y_L(t_0)$$

ie, initially there was a lepton asymmetry yield of $\frac{1}{0.64}$ times the current lepton asymmetry. Therefore;

$$Y_B(t_0) = \frac{-c}{0.64} Y_L(t_0)$$

by putting the value of c in the equation we get the present baryon and lepton asymmetry as

$$Y_B(t_0) = Y_B = -0.553 Y_L \approx -0.55 Y_L$$

which is almost equal to the value we got from the chemical potential and CMB data.

similarly we can write

$$n_B(t_0) = -c n_L(t_i)$$

where c=$\frac{28}{79}$

consider $n_B(t_0)$ as the baryon number density present today, hence we can get the initial lepton number density by using the value

5.0.3 Independent calculations of Baryon to anti-baryon and Lepton to anti-lepton

let us calculate the ratio of baryons to anti-baryons and leptons to anti-leptons from the given data.

consider $n_l = k_1 n_b$ and $n_{\bar{l}} = k_2 n_{\bar{b}}$, which is the fraction of leptons to baryons with a factor k_1 and anti-lepton to anti-baryons with a factor k_2 respectively.

since $\frac{\eta_1}{\eta_2} = \frac{n_B}{n_L} = \frac{(n_b - n_{\bar{b}})}{(n_l - n_{\bar{l}})} = -0.54 \approx -\frac{1}{2}$ substituting for n_L and $n_{\bar{L}}$ in the expression. Hence we have

$$\frac{(n_b - n_{\bar{b}})}{k_1 n_b - k_2 n_{\bar{b}}}) = -\frac{1}{2}$$

by equating the LHS and RHS finally we get

$$2n_b - 2n_{\bar{b}} = (k_2 n_{\bar{b}}) - (k_1 n_b)$$

$$\rightarrow \frac{n_b}{n_{\bar{b}}} = \frac{k_2 + 2}{k_1 + 2}$$

similarly we get $\frac{n_l}{n_{\bar{l}}}$ as

$$\frac{n_l}{n_{\bar{l}}} = \frac{k_1 n_b}{k_2 n_{\bar{b}}} = \frac{k_1}{k_2}\frac{(k_2+2)}{(k_1+2)}$$

we know

$$Y_B(t_0) = \frac{-c}{0.64} Y_L(t_0)$$

and it can be written in the form

$$Y_B(t_0) = -1.56c Y_L(t_0)$$

therefore;

$$\frac{n_b - n_{\bar{b}}}{s} = -1.56c \frac{n_l - n_{\bar{l}}}{s}$$

$$n_b \left(1 - \frac{n_{\bar{b}}}{n_b}\right) = -1.56c \quad n_l \left(1 - \frac{n_{\bar{l}}}{n_l}\right)$$

put the value of ratios of $\frac{n_b}{n_{\bar{b}}}$ and $\frac{n_l}{n_{\bar{l}}}$ respectively

$$n_b \left(1 - \frac{K_1+2}{K_2+2}\right) = -1.56c \quad n_l \left(1 - \frac{k_2(k_1+2)}{k_1(k_2+2)}\right)$$

$$\frac{n_b}{n_l} = \frac{1}{k_1} = -1.56c \quad \frac{\left(1 - \frac{k_2(k_1+2)}{k_1(k_2+2)}\right)}{\left(1 - \frac{K_1+2}{K_2+2}\right)}$$

similarly we get

$$\frac{n_{\bar{b}}}{n_{\bar{l}}} = \frac{1}{k_2} = -1.56c \quad \frac{\left(\frac{k_1(k_2+2)}{k_2(k_1+2)} - 1\right)}{\left(\frac{K_2+2}{K_1+2} - 1\right)}$$

since we know that there are more baryons or leptons than its anti-particles

$$\frac{n_l}{n_b} > \frac{n_{\bar{l}}}{n_{\bar{b}}} \rightarrow k_1 > k_2$$

CHAPTER 6

CONCLUSION

The existence of cosmic matter-antimatter asymmetry is still unclear. The Standard Model of basic particle physics meets all of the essential criteria for baryon asymmetry to emerge from an originally symmetric phase, however the amount of asymmetry that can be produced falls short of what has been seen empirically.

Each of these processes has both appealing and troublesome features, as shown in a short review of GUT scenarios, electroweak baryogenesis, Affleck-Dine and leptogenesis. The addition of a single triplet Higgs to the Standard Model (SM) gives a straightforward framework for explaining inflation, neutrino masses, and baryon asymmetry. The neutral component of the triplet Higgs and SM Higgs cause an inflationary era, with predictions that match data perfectly. This inflationary environment offers the ideal environment for effective Leptogenesis. Because of the finding of neutrino masses, the theory of Leptogenesis has gained fresh traction in recent years. If they are the result of a heavy right neutrino, leptogenesis emerges as a plausible mechanism. The situation is consistent, according to theoretical studies. The degree of arbitrariness is limited in terms of unknown parameters: in order to get the proper order of magnitude of the BAU, the masses of the light neutrinos must fall within a range that is compatible with the interpretation of solar and atmospheric neutrino data. If the light neutrinos turn out to be Majorana particles, the idea gains even more traction.

Whatever the explanation, the universe's baryon asymmetry provides strong cosmological evidence for novel and fascinating phenomena outside standard model of particle.

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